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Thermal Right-Handed Sneutrino Dark Matter in the F_D-Term Model of Hybrid Inflation

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ABSTRACT

We compute the relic abundance of the right-handed sneutrinos in the supersymmetric F_D-term model of hybrid inflation. As well as providing a natural solution to the μ- and gravitino overabundance problems, the F_D-term model offers a new viable candidate to account for the cold dark matter in the Universe: the lightest right-handed sneutrino. In particular, the F_D-term model predicts a new quartic coupling of purely right-handed sneutrinos to the Higgs doublets that thermalizes the sneutrinos and makes them annihilate sufficiently fast to a level compatible with the current cosmic microwave background data. We analyze this scenario in detail and identify favourable regions of the parameter space within the framework of minimal supergravity, for which the lightest right-handed sneutrino becomes the *thermal* dark matter, in agreement with WMAP observations of cosmological inflation. Constraints derived from direct dark matter searches experiments are presented.

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1 Introduction

Hybrid inflation [1], along with its supersymmetric realizations [2–5], remains one of the most predictive and potentially testable scenarios of inflation that have been suggested so far. Hybrid inflation is predictive and testable, in the sense that the inflaton dynamics is mainly governed by a few renormalizable operators which might have observable implications for laboratory experiments. In such a scenario, inflation terminates through the so-called waterfall mechanism, which is triggered, when the inflaton field ϕ passes below some critical value ϕ_c. From that point on, another field X, called the waterfall field, held fixed at origin initially, quickly rolls down to its true vacuum expectation value (VEV) and drastically modifies the slow-roll form of the ϕ-potential, thereby ending inflation.

In supersymmetric theories, the required form of the hybrid inflationary potential may originate from either the F-terms of the superpotential or from a large Fayet–Iliopoulos (FI) D-term [6], usually induced by some anomalous local U(1) symmetry within the context of string theories. In both the F- and D-term hybrid inflation, the slow-roll slope of the potential may come either from supergravity (SUGRA) corrections [2] and/or from radiative effects [3–5].

Recently, a new supersymmetric hybrid inflationary model was proposed in [7] and studied in detail in [8]. The model realizes F-term hybrid inflation and includes a subdominant *non-anomalous* FI D-term that arises from the $\mathrm{U}(1)_X$ gauge symmetry of the waterfall sector. It has therefore been called the F_D-term model of hybrid inflation, or in short, the F_D-term model. The F_D-term model can naturally accommodate the currently favoured red-tilted spectrum with $n_{\rm s} - 1 \approx -0.037$ [9], along with the actual value of the power spectrum of curvature perturbations, $P_{\mathcal{R}} \simeq 4.86 \times 10^{-5}$ [10], and the required number of e-folds, $\mathcal{N}_e \approx 55$ [11].

The presence of the FI term in the F_D-term model is necessary to approximately break a D-parity that governs the waterfall sector. The approximate breaking of the D-parity gives rise to late decays of the superheavy waterfall-sector particles that are produced just after inflation during the preheating epoch [12, 13]. These waterfall particles have masses of the Grand Unified Theory (GUT) scale and can dominate the energy density of the Universe, provided the inflaton coupling κ to the waterfall sector is not too suppressed, i.e. for values of $\kappa \gtrsim 10^{-3}$. Then, the late decays of the GUT-scale waterfall particles produce an enormous entropy that can reduce the gravitino abundance $Y_{\widetilde{G}}$ well below the limits imposed by big bang nucleosynthesis (BBN), i.e. $Y_{\widetilde{G}} \lesssim 10^{-15}$ [14]. In this way, the F_D-term model provides a viable solution to the gravitino overabundance problem [8], without the need to unnaturally suppress all renormalizable inflaton couplings to the particles of

the Minimal Supersymmetric Standard Model (MSSM) sector, below the 10^{-6} level.

Another interesting feature of the F_D-term model is that the μ-parameter of the MSSM can be generated effectively by the superpotential operator $\lambda \widehat{S} \widehat{H}_u \widehat{H}_d$, when the scalar component of the inflaton chiral multiplet $\widehat{S}$ receives a non-zero VEV after the spontaneous symmetry breaking (SSB) of the local $\mathrm{U}(1)_X$ symmetry of the waterfall sector [15]. Moreover, the inflaton superfield $\widehat{S}$ couples to the right-handed neutrino superfields $\widehat{N}_{1,2,3}$, via the superpotential coupling $\frac{1}{2}\rho_{ij}\widehat{S}\widehat{N}_i\widehat{N}_j$, with $i, j = 1, 2, 3$. Hence, the inflaton VEV will produce an effective Majorana mass matrix as well [7, 16]. As a consequence, the resulting heavy Majorana neutrinos are expected to have masses of order μ. If ρ_{ij} is approximately SO(3) symmetric, i.e. $\rho_{ij} \approx \rho\, \mathbf{1}_3$, a possible explanation of the observed baryon asymmetry in the Universe (BAU) may be obtained by thermal electroweak-scale resonant leptogenesis, in a way independent of any pre-existing lepton- or baryon-number asymmetry [17].

Even though the F_D-term model violates explicitly the lepton number (L) by $\Delta L = 2$ superpotential operators, it conserves R-parity. Hence, the lightest supersymmetric particle will be stable and so will potentially qualify as a candidate for the cold dark matter (DM) in the Universe. Most interestingly, the F_D-term model provides a new candidate for the cold DM. This is the lightest right-handed sneutrino (LRHS), which may possess thermal relic abundance [8] for relatively large values of the aforementioned superpotential couplings λ and ρ, i.e. for $\lambda,\, \rho \gtrsim 10^{-2}$. This should be contrasted with what is happening in standard seesaw extensions of the MSSM, where $\widehat{N}_{1,2,3}$ have only bare Majorana masses. Because the small neutrino Yukawa couplings are the only possible interactions of sneutrinos with matter in these models, purely right-handed sneutrinos turn out to be non-thermal and tend to overclose the Universe by many orders of magnitude [18, 19]. It is therefore difficult for the LRHS to be a thermal DM in seesaw extensions of the MSSM with bare Majorana masses.

In this paper we analyze in detail the relic abundance of the right-handed sneutrinos in the supersymmetric F_D-term model of hybrid inflation. In this model, the F-term of the inflaton superfield, F_S, gives rise to the new quartic coupling, $\frac{1}{2}\lambda\rho\, \widetilde{N}_i^* \widetilde{N}_i^* H_u H_d$, in the scalar potential, which involves the right-handed sneutrinos $\widetilde{N}_{1,2,3}$ and the Higgs doublets $H_{u,d}$. As mentioned above, unless the couplings λ and ρ are too small, the new quartic coupling will be sufficiently strong to thermalize the sneutrinos and make them annihilate to a level compatible with the current cosmic microwave background (CMB) data [9], from which the DM component of the Universe was found to be

$$\Omega_{\mathrm{DM}}\, h^2 \;=\; 0.1099 \,\pm\, 0.0062\,. \tag{1.1}$$

The central goal of our analysis is to delineate the parameter space within the context of

minimal supergravity (mSUGRA), for which the LRHS is the *thermal* DM. In addition, we consider the constraints obtained by WMAP observations related to cosmological inflation. Finally, we present numerical estimates of the scattering cross-section of the LRHS with nuclei that will be relevant to direct DM searches in present and future experiments.

After this introduction, the paper is organized as follows: in Section 2 we present the basic structure of the F_D-term model and briefly review the solution to the gravitino overabundance problem. Moreover, in the same section we derive the constraints imposed on the theoretical parameters by cosmological inflation. In Section 3 we perform a detailed study of the relic abundance of the LRHS and offer numerical estimates of representative scenarios within the mSUGRA framework. We also present numerical estimates for the scattering cross-section of the LRHS with the nucleon, indicating the presently achieved and future sensitivity of the current and projected experiments for DM searches, such as CDM-II, SuperCDMS and Xenon1T. Finally, we summarize our conclusions in Section 4.

2 The F_D-Term Model of Hybrid Inflation

In this section we first outline the basic structure of the F_D-term model of hybrid inflation. Then we briefly review how the gravitino abundance can be solved within the F_D-term model. Finally, we present the constraints on the theoretical parameters that are imposed by CMB data pertinent to inflation. A more detailed discussion of all the above issues may be found in [8].

2.1 The Model

The F_D-term model may be defined through the superpotential

$$W \;=\; \kappa\,\widehat{S}\left(\widehat{X}_1\widehat{X}_2\;-\;M^2\right)\;+\;\lambda\,\widehat{S}\widehat{H}_u\widehat{H}_d\;+\;\frac{\rho_{ij}}{2}\,\widehat{S}\,\widehat{N}_i\widehat{N}_j\;+\;h^\nu_{ij}\widehat{L}_i\widehat{H}_u\widehat{N}_j+\;W^{(\mu=0)}_{\rm MSSM}\;, \tag{2.1}$$

where $\widehat{S}$ is the gauge-singlet inflaton superfield and $\widehat{X}_{1,2}$ is a chiral multiplet pair of the so-called waterfall fields which have opposite charges under the $\mathrm{U}(1)_X$ gauge group, i.e. $Q(\widehat{X}_1) = -Q(\widehat{X}_2) = 1$. In addition, $W^{(\mu=0)}_{\rm MSSM}$ indicates the MSSM superpotential without the μ-term,

$$W^{(\mu=0)}_{\rm MSSM}\;=\;h^u_{ij}\,\widehat{Q}_i\widehat{H}_u\widehat{U}_j\;+\;h^d_{ij}\,\widehat{H}_d\widehat{Q}_i\widehat{D}_j\;+\;h_l\,\widehat{H}_d\widehat{L}_l\widehat{E}_l\;. \tag{2.2}$$

Within the SUGRA framework, the sector of soft supersymmetry (SUSY) breaking (SSB) derived from (2.1) is given by

$$-\mathcal{L}_{\rm soft} = M^2_{\tilde{S}} S^* S + M^2_{\tilde{N}} N^*_i N_i + \Big(\kappa A_\kappa\, S X_1 X_2 + \lambda A_\lambda S H_u H_d + \frac{\rho}{2}\, A_\rho\, S \widetilde{N}_i \widetilde{N}_i - \kappa a_S M^2 S + {\rm H.c.}\Big)\,, \tag{2.3}$$

where $M_{\tilde{S}}$, $M_{\tilde{N}}$, $A_{\kappa,\lambda,\rho}$ and a_S are soft SUSY-breaking mass parameters that are all typically of order $M_{\rm SUSY}\ \sim\ 1$ TeV. In addition, the F_D-term model contains a FI D-term, $-\frac{1}{2} g m^2_{\rm FI} D$, associated with the $\mathrm{U}(1)_X$ gauge symmetry of the waterfall sector. The latter gives rise to the D-term potential

$$V_D\ =\ \frac{g^2}{8}\, \Big(|X_1|^2\ -\ |X_2|^2\ -\ m^2_{\rm FI}\Big)^2\ , \tag{2.4}$$

where g is the $\mathrm{U}(1)_X$ gauge-coupling constant. The FI mass parameter $m_{\rm FI}$ is subdominant with respect to the superpotential tadpole mass M, i.e. $m_{\rm FI}/M \lesssim 10^{-5}$.

An interesting feature of the F_D-term model is the generation of an effective μ-term of the required order $M_{\rm SUSY}$ after the SSB of $\mathrm{U}(1)_X$. To see this, let us neglect the VEVs of $H_{u,d}$ next to the large VEVs of the waterfall fields $X_{1,2}$: $\langle X_{1,2}\rangle\ =\ M$. To a good approximation, the VEV of S may then be determined by the following part of the potential:

$$V_S\ =\ |F_{X_1}|^2\ +\ |F_{X_2}|^2\ +\ M^2_S\, S^* S\ +\ \Big[\kappa\, M^2 (A_\kappa - a_S)\, S\ +\ {\rm H.c.}\Big]\ , \tag{2.5}$$

where we have set the waterfall fields $X_{1,2}$ to their actual VEVs. Substituting the F-terms of the waterfall fields,

$$F_{X_{1,2}}\ =\ \kappa S\, \langle X_{2(1)}\rangle\ =\ \kappa M\, S\ , \tag{2.6}$$

into (2.5), we obtain

$$V_S\ =\ \Big(2\kappa^2 M^2\ +\ M^2_S\Big)\, S^* S\ +\ \Big[\kappa\, M^2 (A_\kappa - a_S)\, S\ +\ {\rm H.c.}\Big]\ . \tag{2.7}$$

It is then not difficult to derive from (2.7) that at the present epoch of the Universe, the inflaton field, S, acquires the non-zero VEV

$$\langle S\rangle\ =\ \frac{1}{2\kappa}\, |A_\kappa - a_S|\ +\ \mathcal{O}(M^2_{\rm SUSY}/M)\ , \tag{2.8}$$

in the phase convention that $\langle S\rangle$ is positive. Equation (2.8) implies the effective μ-term

$$\mu\ =\ \lambda\, \langle S\rangle\ \approx\ \frac{\lambda}{2\kappa}\, |A_\kappa - a_S|\ . \tag{2.9}$$

If $\lambda \sim \kappa$, the size of μ-parameter is of order $M_{\rm SUSY}$, as required for a successful electroweak Higgs mechanism.

In addition to the generation of an effective μ-parameter, the third term in (2.1), $\frac{1}{2}\, \rho_{ij}\, \widehat{S}\, \widehat{N}_i \widehat{N}_j$, gives rise to an effective lepton-number-violating Majorana mass matrix,

i.e. $M_N = \rho_{ij}\, v_S$. If we assume that ρ_{ij} is approximately SO(3) symmetric, i.e. $\rho_{ij} \approx \rho\, \mathbf{1}_3$, one obtains 3 nearly degenerate right-handed neutrinos $N_{1,2,3}$, with mass

$$m_N \;=\; \rho\, v_S\ . \tag{2.10}$$

If the couplings λ and ρ are comparable, then the μ-parameter will set the scale for the SO(3)-symmetric Majorana mass m_N, i.e. $m_N \sim \mu$ [7]. Evidently, this will lead to a scenario where the singlet neutrinos $N_{1,2,3}$ have TeV or electroweak-scale masses. This opens up the possibility of directly detecting these singlet Majorana neutrinos through their lepton-number violating signatures at the LHC [20] or ILC [21]. Furthermore, in the F_D-term model the BAU could be explained by thermal electroweak-scale resonant leptogenesis [17].

2.2 Solution to the Gravitino Overabundance Problem

The FI mass term $m_{\rm FI}$ plays a key role in providing a viable solution to the gravitino overabundance problem in the F_D-term model, without the need to unnaturally suppress all the inflaton couplings κ, λ and ρ below the 10^{-6} level [7,8].

In detail, the presence of $m_{\rm FI}$ explicitly breaks an unwanted discrete symmetry that arises from the permutation of the waterfall fields: $\widehat{X}_1 \leftrightarrow \widehat{X}_2$. If $m_{\rm FI}$ was absent, the permutation symmetry would remain exact even after the SSB of the $\mathrm{U}(1)_X$. This would act like parity and was therefore termed D-parity in [7]. As a consequence of D-parity conservation, the D-odd waterfall particles of mass gM would have been stable, and if abundantly produced during the preheating epoch [12,13], they could overclose the Universe at late times.

To avoid this undesirable situation, we introduce a small but non-zero FI term $m_{\rm FI}$. In this case, the D-odd waterfall particles will have forbidden decays to two D-even inflaton-related fields of mass κM, induced by the FI term. To kinematically allow for such decays, we assume that $\kappa < g/2$, where g is the value of the $\mathrm{U}(1)_X$ coupling constant at the GUT scale. The late decays of the D-odd waterfall fields will then reheat again the Universe at temperature T_g, and so release enormous entropy that might be sufficient to reduce the gravitino abundance $Y_{\widetilde{G}}$ below the BBN limits. More explicitly, after the Universe passes through a second reheating phase, the gravitino abundance may be estimated by [8]:

$$Y_{\widetilde{G}} \;\approx\; \frac{7.6\times 10^{-11}}{\kappa g}\,\left(\frac{T_g}{10^{10}\ \mathrm{GeV}}\right)\,, \tag{2.11}$$

Hence, for second reheat temperatures $T_g \sim 1$ TeV and inflaton couplings $\kappa \gtrsim 10^{-2}$, the strict constraint $Y_{\widetilde{G}} \lesssim 10^{-15}$, for $m_{\widetilde{G}} \lesssim 500$ GeV, can be comfortably met.

To determine the second reheat temperature T_g, we may use the standard freeze-out condition $\Gamma_g = H(T_g)$, where

$$\Gamma_g \;=\; \frac{g^4}{128\pi}\,\frac{m_{\rm FI}^4}{M^3} \tag{2.12}$$

is the decay rate of the D-odd particles and

$$H(T) \;=\; \left(\frac{\pi^2 g_*}{90}\right)^{1/2}\frac{T^2}{m_{\rm Pl}} \tag{2.13}$$

is the Hubble expansion parameter in the radiation dominated era of the Universe and $m_{\rm Pl} = 2.4\times 10^{18}$ GeV is the reduced Planck mass. In particular, for a fixed given value of T_g, we may infer the required size of the FI mass term $m_{\rm FI}$ [8]:

$$\frac{m_{\rm FI}}{M} \;\approx\; 8.4\times 10^{-4}\times\left(\frac{0.5}{g}\right)^{3/4}\left(\frac{T_g}{10^9~\text{GeV}}\right)^{1/2}\left(\frac{10^{16}~\text{GeV}}{M}\right)^{1/4}. \tag{2.14}$$

As can be seen from (2.14), for $T_g \sim 1$ TeV, it should be $m_{\rm FI}/M \sim 10^{-6}$, so the FI mass term $m_{\rm FI}$ needs be much smaller than M. Detailed discussion of how such an hierarchy can be naturally achieved within the SUGRA framework may be found in [8].

2.3 Constraints from Cosmological Inflation

Here we recall the constraints derived in [8] on the F_D-term model from cosmological inflation. In fact, there are three constraints that need to be considered.

The first constraint arises from the requirement of solving the horizon and flatness problems of the standard Big-Bang Cosmology. According to the inflationary paradigm, these problems may naturally be solved, if our observable Universe had an accelerated expansion of a number of 50–60 e-folds. In the slow-roll approximation, the number of e-folds, $\mathcal{N}_e$, may be calculated by [11]

$$\mathcal{N}_e \;=\; \frac{1}{m_{\rm Pl}^2}\int_{\phi_{\rm end}}^{\phi_{\rm exit}} d\phi\,\frac{V_{\rm inf}}{V'_{\rm inf}} \;\simeq\; 55\,, \tag{2.15}$$

where $\phi = \sqrt{2}\,{\rm Re}\,S$ is the inflaton field and $V_{\rm inf}$ is the F_D-term inflaton potential that can be found in Section 2.1 of [8]. We will always denote differentiation with respect to ϕ with a prime on $V_{\rm inf}$. Moreover, $\phi_{\rm exit}$ is the value of ϕ, when our present horizon scale exited inflation's horizon, whilst $\phi_{\rm end}$ is its value at the end of inflation. Specifically, the field value $\phi_{\rm end}$ may be determined from the condition:

$$\mathsf{max}\{\epsilon(\phi_{\rm end}), |\eta(\phi_{\rm end})|\} \;=\; 1\,, \tag{2.16}$$

with

$$\epsilon \ = \ \frac{m_{\rm Pl}^2}{2}\left(\frac{V'_{\rm inf}}{V_{\rm inf}}\right)^2 , \qquad \eta \ = \ m_{\rm Pl}^2\,\frac{V''_{\rm inf}}{V_{\rm inf}}\ . \tag{2.17}$$

The other two inflationary constraints come from the so-called power spectrum $P_{\mathcal{R}}$ of curvature perturbations and the spectral index n_s. The square root of the power spectrum, $P_{\mathcal{R}}^{1/2}$, is given by

$$P_{\mathcal{R}}^{1/2} \ = \ \frac{1}{2\sqrt{3}\,\pi m_{\rm Pl}^3}\,\frac{V_{\rm inf}^{3/2}(\phi_{\rm exit})}{|V'_{\rm inf}(\phi_{\rm exit})|}\ . \tag{2.18}$$

This prediction must be compared with the result obtained by a 3-years WMAP analysis of CMB data [10],

$$P_{\mathcal{R}}^{1/2} \ \simeq \ 4.86\times 10^{-5}\,. \tag{2.19}$$

Moreover, in the slow-roll approximation, the spectral index $n_{\rm s}$ is given by [11]

$$n_{\rm s} \ = \ 1 - 6\epsilon(\phi_{\rm exit}) \ + \ 2\eta(\phi_{\rm exit}) \ \simeq \ 1 \ + \ 2\eta(\phi_{\rm exit}), \tag{2.20}$$

where the parameter ϵ is negligible in the F_D-term model. Recently, after analysing its data collected in the last 5 years, WMAP has reported the value for the spectral index [9]:

$$n_{\rm s} \ - \ 1 \ = \ -0.037^{+0.014}_{-0.015}\ . \tag{2.21}$$

This result slightly favours a red-tilted spectrum and is consistent with scale invariance at the 2.64 σ confidence level.

Given the three constraints (2.15), (2.19) and (2.21), and assuming that all inflaton couplings are equal, i.e. $\kappa = \lambda = \rho$, one obtains within mSUGRA the upper bound [8]

$$\kappa \ \lesssim \ 2\times 10^{-2}\ . \tag{2.22}$$

On the other hand, the inflationary scale M is close to the GUT scale, i.e. $M \sim 10^{16}$ GeV, when κ reaches its upper bound imposed by inflation. For an inflaton sector that realizes a next-to-minimal Kähler potential with a negative Hubble-induced mass term for S [22], the upper limit on κ may be slightly relaxed to [8]

$$\kappa \ \lesssim \ 3.2\times 10^{-2}\ , \tag{2.23}$$

whilst M decreases to $M \simeq 0.5\times 10^{16}$ GeV.

It is important to properly translate the upper bounds (2.22) and (2.23) on κ obtained at the inflationary scale M into the respective ones on λ and ρ for the soft SUSY-breaking scale $M_{\rm SUSY}$. As we will see more explicitly in the next section, it is the product $\lambda\rho$ evaluated at the scale $M_{\rm SUSY}$ that controls the strength of annihilation of the LRHSs into

the Higgs fields and other SM particles. Even though the renormalization group (RG) evolution of ρ from M to $M_{\rm SUSY}$ may be ignored, as $\rho(M) \approx \rho(M_{\rm SUSY})$, this is not the case for the coupling λ. Neglecting gauge and small Yukawa couplings of order 10^{-1}, the RG equation for λ is given by [23]

$$16\pi^2 \frac{d\lambda}{dt} \;=\; \lambda\left(\frac{3}{2}\,h_t^2 \;+\; \frac{3}{2}\,h_b^2\right)\,, \tag{2.24}$$

where $t = \ln(Q^2/M_{\rm SUSY})$. Assuming that the RG evolution is dominated by the top-quark Yukawa coupling h_t, the solution to (2.24) is easily found to be

$$\lambda(M_{\rm SUSY}) \;=\; \lambda(M)\left(\frac{M_{\rm SUSY}}{M}\right)^{3h_t^2/(16\pi^2)} \;\approx\; 0.57\times\lambda(M)\;. \tag{2.25}$$

To obtain the last result in (2.25), we assumed that $h_t \approx 1$ and $M_{\rm SUSY}/M \sim 10^{-13}$. Then, starting with the boundary condition $\lambda = \kappa$ at the inflationary scale M, the RG running (2.25) of λ implies the upper limits:

$$\lambda(M_{\rm SUSY}) \;\lesssim\; 1.14\times 10^{-2}\,, \qquad \lambda(M_{\rm SUSY}) \;\lesssim\; 1.82\times 10^{-2}\;, \tag{2.26}$$

for an inflaton sector with a minimal and a next-to-minimal Kähler potential, respectively.

In addition to constraints from inflation, one may also get constraints on the size of M from cosmic strings that arise due to the SSB of the local $\mathrm{U}(1)_X$ symmetry. For values of $\kappa \sim 10^{-2}$ of our interest, this implies that one must have [8] $M \lesssim 0.5\times 10^{16}$ GeV. This constraint may be a bit restrictive for the mSUGRA model, but it can be completely avoided if the waterfall sector realizes an $\mathrm{SU}(2)_X$ gauge symmetry instead of $\mathrm{U}(1)_X$, whose SSB generates no topological defects [8]. Consequently, we will conservatively consider the limits stated in (2.22), (2.23) and (2.26) when implementing inflationary constraints on the relic abundance of the LRHS in the next section.

3 Right-Handed Sneutrino as Thermal Dark Matter

In the F_D-term hybrid model R-parity is conserved, even though the lepton number L, as well as $B-L$, are explicitly broken by the Majorana operator $\frac{1}{2}\rho\widehat{S}\widehat{N}\widehat{N}$. We note that all superpotential couplings either conserve the $B-L$ number or break it by even number of units. Since R-parity of each superpotential operator is determined to be $R = (-1)^{3(B-L)} = +1$, the F_D-term hybrid model conserves R-parity. As a consequence, the LSP of the spectrum is stable and can be a viable candidate for Cold Dark Matter (CDM). As an

extension of the MSSM, our model can accommodate the standard SUSY CDM candidates, such as the lightest neutralino. Because of the connection between the Higgs and neutrino sectors, on the one hand, and inflation, on the other, it is very interesting to explore the possibility of having a right-handed sneutrino as LSP in order to solve the CDM problem. As we will see in Section 3.3, this renders the F_D-term model much more constrained, leading to sharp predictions for scattering cross-sections relevant to experiments of direct searches for CDM.

3.1 Sneutrino Mass Spectrum

Before calculating the sneutrino relic abundance in our model, we first observe that light right-handed sneutrinos may easily appear in the spectrum. Ignoring the terms proportional to the small neutrino-Yukawa couplings, the 6×6 right-handed sneutrino mass matrix $\mathcal{M}^2_{\widetilde{N}}$ is given in the weak basis $(\widetilde{N}_{1,2,3}, \widetilde{N}^*_{1,2,3})$ by

$$\mathcal{M}^2_{\widetilde{N}} = \frac{1}{2} \begin{pmatrix} \rho^2 v_S^2 + M^2_{\widetilde{N}} & \rho A_\rho v_S + \rho\lambda v_u v_d \\ \rho A^*_\rho v_S + \rho\lambda v_u v_d & \rho^2 v_S^2 + M^2_{\widetilde{N}} \end{pmatrix}, \tag{3.1}$$

where $v_S = \langle S \rangle$, $v_{u,d} = \langle H_{u,d} \rangle$. Moreover, $M^2_{\widetilde{N}}$ is the soft SUSY-breaking mass matrix associated with the sneutrino fields and A_ρ is the sneutrino trilinear coupling matrix. In general, $\mathcal{M}^2_{\widetilde{N}}$ is diagonalized by a unitary matrix $U_{\widetilde{N}}$ such that

$$U^\dagger_{\widetilde{N}} \mathcal{M}^2_{\widetilde{N}} U_{\widetilde{N}} = \mathrm{diag}\left(m^2_{\tilde{N}_1}, m^2_{\tilde{N}_2}, \ldots, m^2_{\tilde{N}_6} \right), \tag{3.2}$$

where the sneutrino masses are ordered, such that $m_{\tilde{N}_1} < m_{\tilde{N}_2} < \ldots < m_{\tilde{N}_6}$. Neglecting the possible flavor structure contained in the 3×3 matrices $M^2_{\widetilde{N}}$ and A_ρ, the sneutrino spectrum will then consist of 3 light (heavy) right-handed sneutrinos with masses

$$m^2_{\tilde{N}_{L(H)}} = \rho^2 v_S^2 + M^2_{\widetilde{N}} - (+) \left| \rho A_\rho v_S + \rho\lambda v_u v_d \right|. \tag{3.3}$$

All mass terms in (3.3) are $\mathcal{O}(100\text{–}1000)$ GeV, so a proper choice of model parameters can accommodate a LRHS to act as LSP. Unless the trilinear coupling A_ρ is small compared to μ, the off-diagonal elements in (3.1) will induce a sizeable mixing between the heavy and light right-handed sneutrino states, suppressing the light masses to values smaller than $(\mu^2 + M^2_{\widetilde{N}})^{1/2}$. This will be demonstrated in our discussion of the numerical results in Section 3.3, where the F_D-term model is embedded within the mSUGRA framework.

3.2 Sneutrino Annihilation and Relic Density

Right-handed sneutrinos as CDM were considered in [18] in the context of the MSSM with right-handed neutrino superfields $\widehat{N}_i$ and bare Majorana masses $M^N_{ij} \widehat{N}_i \widehat{N}_j$. This analysis

shows that thermal right-handed sneutrinos have rather high relic abundances and will generally overclose the Universe. The reason is that because of the small neutrino Yukawa couplings h^ν_{ij}, the self- and co-annihilation interactions of the sneutrino LSP with itself and other MSSM particles are rather weak. These weak processes do not allow the sneutrino LSP to stay long enough in thermal equilibrium before its freeze-out temperature, such that its number density gets reduced to the observed value $\Omega_{\rm DM}h^2 \approx 0.11$ [9] [cf. (1.1)]. In fact, the predicted values for $\Omega_{\rm DM}h^2$ turn out to be many orders of magnitude larger than 1. Instead, right-handed sneutrinos can be viable thermal DM candidates in the MSSM if they significantly mix with left-handed sneutrinos, either by increasing the SUSY-breaking trilinear couplings [24] [1], or by lowering the right-handed neutrino mass scale [25]. Alternatively, right-handed sneutrinos may become thermal DM by introducing a new U(1)' gauge coupling to make the self-annihilation interaction sufficiently strong [19]. Recently, there has been a paper [26] discussing the possibility of right-handed sneutrinos as DM in an extended version of the next-to-minimal supersymmetric standard model.

In the F_D-term hybrid model, a novel possibility opens up. As was first observed in [8], there exists a new quartic coupling described by the Lagrangian [2]

$$\mathcal{L}^{\rm LSP}_{\rm int} = \frac{1}{2}\lambda\rho\tilde{N}^*_i\tilde{N}^*_i H_u H_d + {\rm H.c.}\ . \tag{3.4}$$

This quartic coupling between right-handed sneutrinos and Higgs fields results from the F-term of the inflaton field F_S: $\frac{1}{2}\rho\widehat{N}_i\widehat{N}_i + \lambda\widehat{H}_u\widehat{H}_d \subset F_S$. If strong enough, the interaction (3.4) can thermalize the sneutrinos and make them annihilate to a level compatible with the current CMB data via the processes depicted in Figure 1.

For sneutrino masses of our interest, the most relevant processes are the off-resonant pair-production of W bosons and the on-shell pair-production of light Higgs bosons. An initial estimate of the process $\tilde{N}\tilde{N} \to \langle H_u\rangle H_d \to W^+W^-$ for $m_{\tilde{N}} > m_W$ yields

$$\Omega_{\rm DM}\, h^2 \ \approx\ \left(\frac{10^{-4}}{\rho^2\lambda^2}\right)\left(\frac{\tan\beta\, m_H}{g_W\, m_W}\right)^2\ . \tag{3.5}$$

In order to obtain an acceptable CDM density, relatively large couplings ρ and λ are needed, $\rho\lambda \gtrsim 0.1$. However, these large values for λ and ρ are not compatible with the constraints derived by inflation.

The situation differs for sneutrino masses $m_{\tilde{N}} < m_W$, in large $\tan\beta$ scenarios, in which light Higgs bosons couple appreciably to b-quarks [36]. In particular, in the kinematic region

[1] For an earlier discussion, see also the paper by N. Arkani-Hamed *et al.* in [16].

[2] The implications of a generic singlet-Higgs quartic coupling for the CDM abundance and detection were studied before in [27–29], within a simple non-SUSY model.

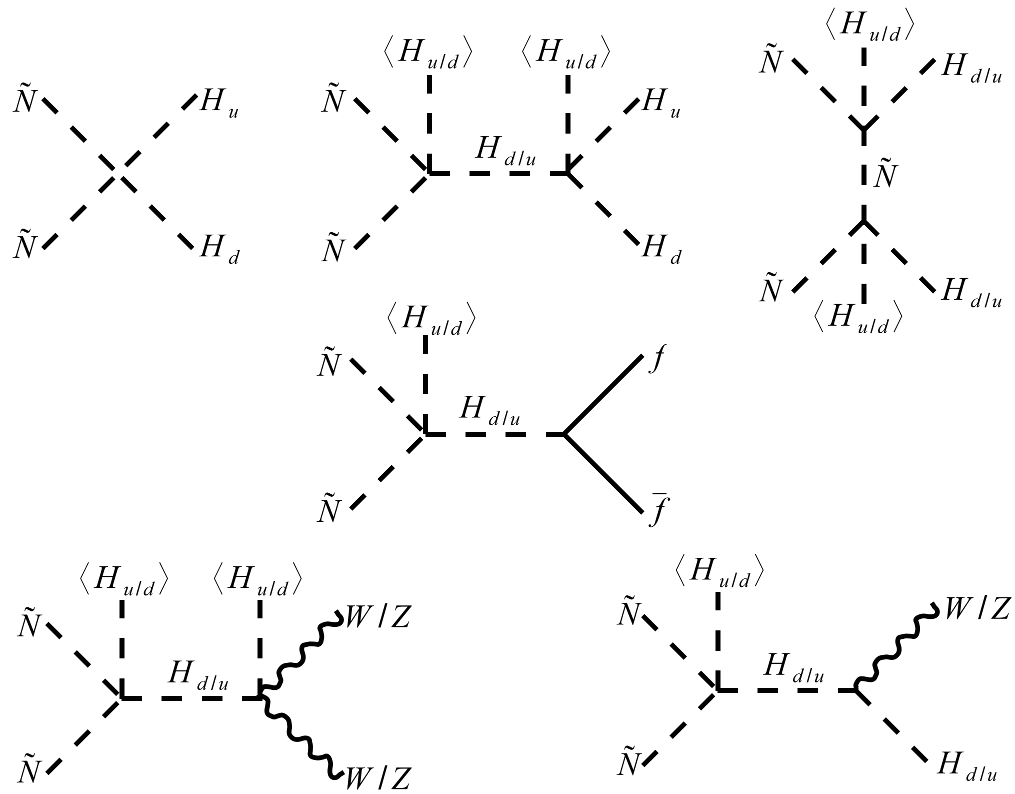

Figure 1: Feynman graphs related to sneutrino annihilation.

$m_{H_1} \approx 2m_{\widetilde{N}_1}$, the self-annihilation process $\widetilde{N}_1\widetilde{N}_1 \to \langle H_u \rangle H_d \to b\bar{b}$ becomes resonant, and the above estimate modifies to

$$\Omega_{\rm DM}h^2 \approx 10^{-4} \times B^{-1}(H_1 \to \widetilde{N}_1\widetilde{N}_1) \times \left(\frac{m_{H_1}}{100\ \text{GeV}}\right)^2 . \tag{3.6}$$

Consequently, if the couplings λ, ρ are not too small, e.g. $\lambda\rho \gtrsim 10^{-3}$, the right-handed sneutrino $\widetilde{N}_1$ can now efficiently annihilate via a resonant H_1-boson into pairs of b-quarks, thus obtaining a relic DM density compatible with the observed value (1.1).

We will now show that the naive estimates (3.5) and (3.6) presented in [8] are in a fairly good agreement with a complete calculation of all relevant sneutrino annihilation processes displayed in Figure 1. To this end, we use the short-hand notation $M_{XY} = M(\tilde{N}_a\tilde{N}_b \to XY)$ to denote the individual matrix elements for the annihilation of sneutrinos $\tilde{N}_a$ and $\tilde{N}_b$. The contributing processes may be listed as follows ($c_w = \cos\theta_w$, $v = 2m_W/g_w$):

(i) $\tilde{N}_a\tilde{N}_b \longrightarrow H^+H^-$, via contact quartic interaction and s-channel Higgs exchange:

$$M_{H^+H^-} \;=\; g_{\tilde{N}_a\tilde{N}_bH^+H^-} - v^2 \sum_{k=1}^{3} \frac{g_{\tilde{N}_a\tilde{N}_bH_k}g_{H_kH^+H^-}}{s - m_{H_k}^2 + im_{H_k}\Gamma_{H_k}} \;; \tag{3.7}$$

(ii) $\tilde{N}_a\tilde{N}_b \longrightarrow W^+W^-$, via s-channel Higgs exchange:

$$M_{W^+W^-} \;=\; g_w m_W v \left[2 \;+\; \left(1-\frac{s}{2m_W^2}\right)^2\right]^{1/2} \sum_{k=1}^{3} \frac{g_{\tilde{N}_a\tilde{N}_b H_k} g_{H_k VV}}{s-m_{H_k}^2+im_{H_k}\Gamma_{H_k}}\;; \qquad (3.8)$$

(iii) $\tilde{N}_a\tilde{N}_b \longrightarrow ZZ$, via s-channel Higgs exchange:

$$M_{ZZ} \;=\; \frac{g_w m_W v}{2c_w}\left[2 \;+\; \left(1-\frac{s}{2m_Z^2}\right)^2\right]^{1/2} \sum_{k=1}^{3} \frac{g_{\tilde{N}_a\tilde{N}_b H_k} g_{H_k VV}}{s-m_{H_k}^2+im_{H_k}\Gamma_{H_k}}\;; \qquad (3.9)$$

(iv) $\tilde{N}_a\tilde{N}_b \longrightarrow f\bar{f}$, via s-channel Higgs exchange:

$$M_{f_\alpha\bar{f}_\alpha} \;=\; v\sqrt{2s}\left[|A_S|^2\left(1-\frac{4m_\alpha^2}{s}\right) \;+\; |A_P|^2\right]^{1/2}, \qquad (3.10)$$

with

$$A_{S/P} \;=\; \sum_{k=1}^{3} \frac{g_{\tilde{N}_a\tilde{N}_b H_k} g_{f_\alpha} g^{S/P}_{H_k\bar{f}_\alpha f_\alpha}}{s-m_{H_k}^2+im_{H_k}\Gamma_{H_k}}\;; \qquad (3.11)$$

(v) $\tilde{N}_a\tilde{N}_b \longrightarrow H_iH_j$, via contact quartic interaction, s-channel Higgs exchange and t/u-channel sneutrino exchange:

$$\begin{aligned} M_{H_iH_j} \;=\;& g_{N_aN_bH_iH_j} \;-\; v^2 \sum_{k=1}^{3} \frac{g_{\tilde{N}_a\tilde{N}_b H_k} g_{H_iH_jH_k}}{s-m_{H_k}^2+im_{H_k}\Gamma_{H_k}} \\ &-\; v^2 \sum_{c=1}^{6} \frac{g_{\tilde{N}_a\tilde{N}_c H_i} g_{\tilde{N}_b\tilde{N}_c H_j}}{t-m_{\tilde{N}_c}^2} \;-\; v^2 \sum_{c=1}^{6} \frac{g_{\tilde{N}_a\tilde{N}_c H_j} g_{\tilde{N}_b\tilde{N}_c H_i}}{u-m_{\tilde{N}_c}^2}\;; \end{aligned} \qquad (3.12)$$

(vi) $\tilde{N}_a\tilde{N}_b \longrightarrow H^+W^-$, via s-channel Higgs exchange:

$$M_{H^+W^-} \;=\; \frac{g_w v}{2}\left[\frac{s^2}{4m_W^2}\left(1-\frac{m_W^2+m_{H^+}^2}{s}\right)^2 \;-\; m_{H^+}^2\right]^{1/2} \sum_{k=1}^{3} \frac{g_{\tilde{N}_a\tilde{N}_b H_k} g_{H_k H^+ W^-}}{s-m_{H_k}^2+im_{H_k}\Gamma_{H_k}}\;; \qquad (3.13)$$

(vii) $\tilde{N}_a\tilde{N}_b \longrightarrow H_iZ$, via s-channel Higgs exchange:

$$M_{H_iZ} \;=\; \frac{g_w v}{4c_s}\left[\frac{s^2}{4m_Z^2}\left(1-\frac{m_Z^2+m_{H_i}^2}{s}\right)^2 - m_{H_i}^2\right]^{1/2} \sum_{k=1}^{3} \frac{g_{\tilde{N}_a\tilde{N}_b H_k} g_{H_k H_i Z}}{s-m_{H_k}^2+im_{H_k}\Gamma_{H_k}}\,. \qquad (3.14)$$

In the above, the effective sneutrino-to-Higgs couplings $g_{\tilde N_a \tilde N_b H^+ H^-}$, $g_{\tilde N_a \tilde N_b H_j H_j}$ and $g_{\tilde N_a \tilde N_c H_i}$ that arise from the interaction Lagrangian (3.4) are given by ($c_\beta = \cos\beta, s_\beta = \sin\beta$)

$$
\begin{aligned}
g_{\tilde N_a \tilde N_b H^+ H^-} &= \frac{\lambda\rho}{2}\, c_\beta s_\beta \delta_{ab}, && (3.15)\\
g_{\tilde N_a \tilde N_b H_i H_j} &= \frac{\lambda\rho}{2}\,\frac{\delta_{ab}}{1+\delta_{ij}}\,[(O_{\phi_u i}O_{\phi_d j} + O_{ai}O_{\phi_u j}s_\beta + O_{ai}O_{\phi_d j}c_\beta - O_{ai}O_{aj}s_\beta c_\beta) \\
&\qquad + (i \leftrightarrow j)]\,, && (3.16)\\
g_{\tilde N_a \tilde N_b H_i} &= \frac{\lambda\rho}{2}\,(O_{\phi_d i}s_\beta + O_{\phi_u i}c_\beta)\,\delta_{ab}\,, && (3.17)
\end{aligned}
$$

where O is the 3×3 Higgs-boson mixing matrix, defined such that

$$
(\phi_d,\, \phi_u,\, a)^T \;=\; O\,(H_1,\, H_2,\, H_3)^T\,. \qquad (3.18)
$$

For the effective Higgs-boson couplings $g_{H_k H_i Z}$, $g_{H_k H^+ W^-}$, $g_{H_i H_j H_k}$, $g_{H_k H^+ H^-}$, $g_{H_k VV}$ and $g_{f_\alpha}\, g^{S/P}_{H_k f_\alpha \bar f_\alpha}$, including O, the Higgs-boson masses $m_{H_{1,2,3}}$ and their decay widths $\Gamma_{H_{1,2,3}}$, we follow the notations and conventions of [30, 31] and calculate them by means of the computational package `CPsuperH`.

The total annihilation cross-section $\sigma_{ab} = \sigma(\tilde N_a \tilde N_b \to \text{all})$ may then be conveniently expressed as the sum of all channels,

$$
\sigma_{ab} \;=\; \sigma_{H^+H^-} + \sigma_{W^+W^-} + \sigma_{ZZ} + \sigma_{H^+W^-} + \sigma_{H^-W^+} + \sum_{i=1}^{3}\sigma_{H_i Z} + \sum_{i,j=1}^{3}\sigma_{H_i H_j} + \sum_{f=\tau,b,t}\sigma_{f\bar f}\,. \qquad (3.19)
$$

The individual cross sections σ_{XY} are defined by

$$
\sigma_{XY} \;=\; \frac{1}{1+\delta_{XY}}\,\frac{1}{16\pi\lambda(s, m^2_{\tilde N_a}, m^2_{\tilde N_b})}\int_{t^-}^{t^+} dt\,|M_{XY}|^2, \qquad (3.20)
$$

with

$$
\begin{aligned}
t^{\pm} &= m_X^2 + m^2_{\tilde N_a} - \frac{1}{2s}\Big((s + m^2_{\tilde N_a} - m^2_{\tilde N_b})(s + m_X^2 - m_Y^2) \\
&\qquad \mp \lambda^{1/2}(s, m^2_{\tilde N_a}, m^2_{\tilde N_b})\lambda^{1/2}(s, m_X^2, m_Y^2)\Big)\,, && (3.21)\\
\lambda(a,b,c) &= (a-b-c)^2 - 4bc. && (3.22)
\end{aligned}
$$

In order to calculate the relic density, we follow [32] and use an effective cross-section averaged over all initial sneutrino channels,

$$
\sigma_{\text{eff}} \;=\; \sum_{a,b=1}^{6}\sigma_{ab}\frac{g_a g_b}{g^2_{\text{eff}}}(1+\Delta_a)^{3/2}(1+\Delta_b)^{3/2}\;\exp\,[-x(\Delta_a+\Delta_b)]\,, \qquad (3.23)
$$

where

$$g_{\rm eff} \;=\; \sum_{a=1}^{6} g_a (1+\Delta_a)^{3/2} e^{-x\Delta_a}\;, \qquad \Delta_a \;=\; \frac{m_{\tilde{N}_a} - m_{\tilde{N}_1}}{m_{\tilde{N}_1}}\;. \tag{3.24}$$

In (3.23), both the effects of LSP self-annihilation and co-annihilation with the heavier sneutrinos are included[3]. In terms of the effective cross-section (3.23), the thermally-averaged effective cross-section may be calculated as

$$\langle \sigma v \rangle \;=\; \frac{x^{3/2}}{2\pi^{3/2}} \int_0^\infty dv\, v^2 (\sigma_{\rm eff} v)\, e^{-xv^2/4}\;, \tag{3.25}$$

where the integrand is expressed in terms of the relative velocity v, such that

$$s \;=\; \frac{4m^2_{\tilde{N}_1}}{1 - v^2/4}\;. \tag{3.26}$$

From the expression (3.25), we may determine the freeze-out temperature $x_f = m_{\tilde{N}_1}/T_f$ by iteratively solving the equation

$$x_f \;=\; \ln\left(\frac{0.038 g_{\rm eff}\, M_{\rm Pl}\, m_{\tilde{N}_1} \langle \sigma v\rangle}{g_*^{1/2} x_f^{1/2}}\right)\;, \tag{3.27}$$

where $M_{Pl} = 1.22 \times 10^{19}$ GeV is the Planck mass and g_* is the total number of effective relativistic degrees of freedom at the temperature of the LSP freeze-out. The present day relic density is then given by

$$\Omega_{\rm DM}\, h^2 \;\approx\; \frac{1.07 \times 10^9 \ {\rm GeV}^{-1}}{J\, g_*^{1/2} m_{Pl}} \tag{3.28}$$

where J is the post freeze-out annihilation efficiency factor given by

$$J \;=\; \int_0^\infty dv\, v (\sigma_{eff} v)\, {\rm erfc}(v\sqrt{x_f}/2). \tag{3.29}$$

In our numerical estimates, we neglect the flavor structure of the right-handed sneutrinos and treat the three light right-handed sneutrinos $\tilde{N}_{1,2,3}$ as being essentially degenerate [4]. Since all three light sneutrinos will contribute to the relic density, we must therefore multiply (3.28) by 3 to obtain the final relic DM abundance.

---

[3]Note that co-annihilation effects become significant, only if the mass differences with the heavier sneutrinos are smaller or comparable to the LSP freeze-out temperature, i.e, when $m_{\tilde{N}_a} - m_{\tilde{N}_b} \lesssim T_f$.

[4]Note that the second and third right-handed sneutrinos $\tilde{N}_{2,3}$ will decay to the LRHS $\tilde{N}_1$ through the processes $\tilde{N}_{2,3} \to \tilde{N}_1 \gamma, \tilde{N}_1 \nu\bar{\nu}$. We do not address potential problems for BBN from the late decays of $\tilde{N}_{2,3}$, since their rates strongly depend on the flavor structure of ρ_{ij} and the Yukawa couplings h^ν_{ij} [cf. (2.1)] and on the details of the model in general.

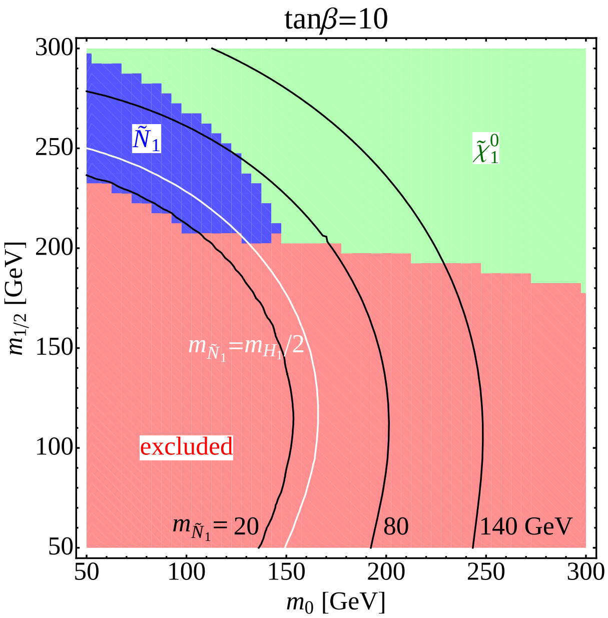

Figure 2: *Allowed* $(m_0,\, m_{1/2})$ *parameter space for a mSUGRA scenario with* $A_0 = 300$ *GeV, sign*$\mu = +$, $\lambda = \rho = 10^{-2}$ *and* $\tan\beta = 10$ *(left panel) and 30 (right panel). The black contours show the predicted LRHS mass, while the sneutrino* $\tilde{N}_1$/*neutralino* $\tilde{\chi}_1^0$/*stau* $\tilde{\tau}_1$ *LSP is given by the blue/green/orange area. The red area is excluded by direct SUSY mass searches. The white contour is defined by the condition* $m_{\tilde{N}_1} = m_{H_1}/2$, *allowing for rapid sneutrino annihilation via the* H_1*-boson resonance.*

3.3 Numerical Results

The numerical analysis is separated in two parts: in the first part, we perform a scan over the mSUGRA parameter space to calculate the supersymmetric particle spectrum and identify regions where the LRHS can be a possible candidate for CDM. In the second part, we specify two mSUGRA scenarios and calculate the constraints on the effective sneutrino annihilation coupling $\lambda\rho$ by requiring a sneutrino relic density of $\Omega_{\rm DM}h^2 = 0.11$.

In Figure 2 we plot the lightest sneutrino mass $m_{\tilde{N}_1}$ as contours in the mSUGRA parameter plane $(m_0, m_{1/2})$, for two different values of $\tan\beta = 10$ (left) and 30 (right). In both plots of Figure 2, we set $A_0 = 300$ GeV and $\mu > 0$. For the inflaton couplings λ, ρ required to calculate the sneutrino masses (3.3), we simply choose

$$\lambda = \rho = 10^{-2}, \tag{3.30}$$

in accordance with the bounds (2.26) derived from inflation. The coloured areas in Figure 2 denote the LSP in the given parameter region: sneutrino $\tilde{N}_1$ (blue), neutralino $\tilde{\chi}_1^0$ (green) or stau $\tilde{\tau}_1$ (orange). The red area on the bottom/left is excluded by direct searches for

SUSY particles. Specifically, the following experimental mass limits are used [34]:

$$\begin{aligned} m_{\tilde{\chi}_1^-} &> 104 \text{ GeV} , \\ m_{\tilde{q}} &> 375 \text{ GeV} , \\ m_{\tilde{g}} &> 289 \text{ GeV} , \\ m_{\tilde{\ell}} &> 95 \text{ GeV} , \\ m_{\tilde{\nu}_L} &> 130 \text{ GeV} . \end{aligned} \tag{3.31}$$

Figure 2 was determined by appropriately using universal soft SUSY-breaking parameters at the GUT scale according to the mSUGRA scheme and then solving the MSSM RG equations down to the electroweak scale. In this respect, our computation was aided by the software package SPheno [35]. We neglect the RG running of the sneutrino parameters $M^2_{\tilde{N}}$ and A_ρ which enter the sneutrino mass matrix (3.1), and identify them directly with m_0^2 and A_0, respectively. This is a reasonable approximation as their RG evolution is only driven by the small couplings λ and ρ. The Higgs coupling parameter μ is then calculated consistently by requiring proper electroweak symmetry breaking. In the F_D-term model, the μ term originates from the VEV of the inflaton (2.9). This immediately allows us to calculate both the inflaton VEV, $\langle S\rangle = \mu/\lambda$, and the mass scale of the right-handed neutrinos, $m_N = \rho\langle S\rangle = \frac{\rho}{\lambda}\mu$ (2.1). For the $\tilde{N}_1$ LSP region of interest and with our choice $\lambda = \rho = 10^{-2}$, μ and m_N are equal and of order 300 GeV. The mass $m_{\tilde{N}_1}$ of the LRHS as LSP ranges between 20–100 GeV. This allows for a rapid annihilation of $\tilde{N}_1$ via the Higgs resonance, $m_{\tilde{N}_1} = m_{H_1}/2 \approx 57$ GeV, along the white contour in Figure 2.

The F_D-term model puts strong constraints on the mSUGRA parameter space, when requiring a sneutrino LSP and taking into account bounds from inflation. As can be seen in Figure 2, the connection between LRHS mass $\tilde{N}_1$ and μ generally points towards a low-energy SUSY spectrum. This coincidentally includes the H_1-boson funnel region, where $m_{H_1} \approx 2m_{\tilde{N}_1}$. On the other hand, very large and small values for A_0 and $\tan\beta$ are disfavoured as they generally exclude a sneutrino LSP. The above correlations may be somewhat relaxed if non-universal inflaton couplings λ and ρ are considered.

In order to compute the sneutrino relic density and analyze the constraints on the effective annihilation coupling $\lambda\rho$, the following two mSUGRA scenarios have been selected:

- Scenario I:

$$m_0 = 70 \text{ GeV},\ m_{1/2} = 243 \text{ GeV},\ A_0 = 300 \text{ GeV},\ \tan\beta = 10,\ \mu = 303 \text{ GeV} . \tag{3.32}$$

- Scenario II:

$$m_0 = 125 \text{ GeV},\ m_{1/2} = 212 \text{ GeV}, A_0 = 300 \text{ GeV},\ \tan\beta = 30,\ \mu = 263 \text{ GeV} . \tag{3.33}$$

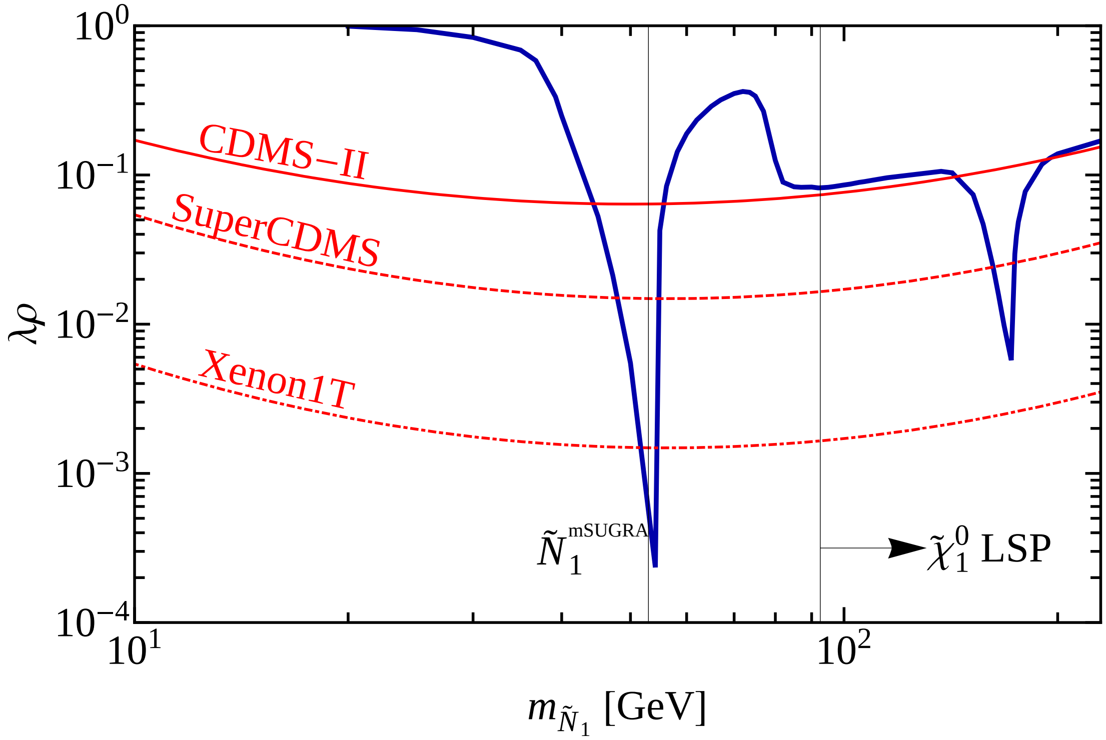

Figure 3: *Effective annihilation coupling $\lambda\rho$ as a function of the mass of the LRHS $m_{\tilde{N}_1}$ for the observed relic density $\Omega_{\rm DM}h^2 = 0.11$ (blue curve) in the mSUGRA Scenario I (3.32). The actual sneutrino and neutralino masses in the scenario are indicated by vertical lines. The red curves denote the upper bound on $\lambda\rho$ as obtained by the CDMS-II experiment and as expected by the projected sensitivities of SuperCDMS and Xenon1T.*

In addition, we keep the LRHS mass as a free parameter. The effective annihilation coupling $\lambda\rho$ (3.4) is then consistently calculated so as to obtain a sneutrino relic density $\Omega_{\rm DM}\, h^2 = 0.11$, consistent with observation. Furthermore, we assume that the mass splitting between the light and heavy right-handed sneutrinos is sufficiently large so that co-annihilation can be safely ignored. This is valid as long as there is a sizeable mixing between the light and heavy right-handed sneutrino states, which is certainly true for the mass range $m_{\tilde{N}_1} < m_{\tilde{\chi}_1^0}$ of our interest. All other MSSM parameters and masses were calculated within the mSUGRA framework. Numerical estimates of the allowed parameters in the $(m_{\tilde{N}_1}, \lambda\rho)$-plane are shown for Scenarios I and II in Figures 3 and Figure 4, respectively.

As we have seen in Section 2.3, the requirement for successful inflation puts upper limits on the couplings λ and ρ. Given (2.22), (2.23) and 2.26), the upper limits on the product $\lambda\rho$ for an inflaton sector with a minimal and next-to-minimal Kähler potential

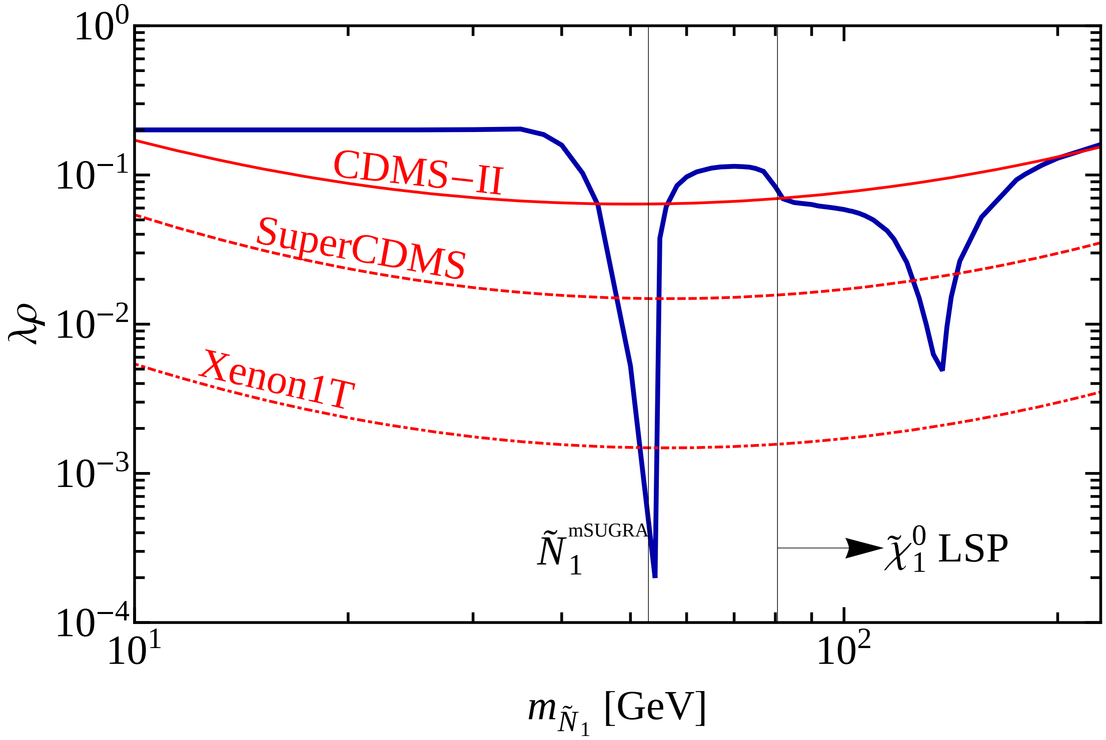

Figure 4: *As in Figure 3, but for the mSUGRA Scenario II (3.33).*

may easily be deduced to be

$$\lambda\rho \ \lesssim\ 2.3\times 10^{-4}\,, \qquad \lambda\rho \ \lesssim\ 5.8\times 10^{-4}\,, \tag{3.34}$$

respectively, at the soft SUSY-breaking scale $M_{\rm SUSY}$. On the other hand, Figures 3 and 4 show that it should be

$$\lambda\rho \ \gtrsim\ 2\times 10^{-4}\,, \tag{3.35}$$

in order to account for the observed DM relic abundance in the H_1-boson funnel region, where $m_{\tilde{N}_1} \approx m_{H_1}/2$. Larger values of $\tan\beta$ do suppress the coupling required to get the observed relic density, but not to a level compatible with the inflationary constraints (3.34). In general, we find that LRHS masses larger than about 100 GeV are not possible within a mSUGRA realization of the F_D-term model. This is indicated by the value of the neutralino mass in the given mSUGRA scenario as displayed by vertical lines in Figures 3 and 4.

Further constraints on the $(m_{\tilde{N}_1}, \lambda\rho)$-plane may be obtained by taking into account the limits from direct searches of experiments which look for scattering between Weakly Interacting Massive Particles (WIMPs) and nuclei. Specifically, a WIMP, such as the LRHS, can directly be detected through its elastic scattering with a nucleus. In our case,

the relevant scattering process is $\tilde{N}_1 + {}^A_Z X \rightarrow \tilde{N}_1 + {}^A_Z X$ and proceeds via a Higgs-boson t-channel exchange. Its cross-section may well be estimated by [29]

$$\sigma_{\rm el}^{\rm nucleus} \;\approx\; \frac{(1/2\lambda\rho)^2 v^2 |M_X|^2}{\pi} \, \frac{m_{\rm red}^2}{m_{\tilde{N}_1}^2 m_{H_1}^4} \,, \tag{3.36}$$

where $m_{\rm red}$ is the reduced mass of the LRHS-nucleus system, i.e.

$$m_{\rm red} \;=\; \frac{m_{\tilde{N}_1} m_X}{m_{\tilde{N}_1} + m_X}, \tag{3.37}$$

and M_X is the nuclear matrix element. For comparison purposes, we express our results in terms of the *nucleon* cross section. Assuming the nucleus to be composed of A independent nucleons, the nuclear cross sections then simply scale quadratically with the nucleon number A and the reduced masses: $m_{\rm red}^2(p)\sigma_{\rm el}^{\rm nucleus} = A^2 m_{\rm red}^2({}^A_Z X)\sigma_{\rm el}^{\rm nucleon}$. The nucleon matrix element $M_{\rm nucleon} \sim 10^{-3}$ is mostly sensitive to the strange-quark Yukawa coupling. An adequate estimate of the elastic scattering cross section $\sigma_{\rm el}^{\rm nucleon}$ of a right-handed sneutrino with a nucleon yields [29]

$$\sigma_{\rm el}^{\rm nucleon} \;\approx\; \left(5\times 10^{-50}\ \ {\rm cm}^2\right) \left(\frac{\lambda\rho}{10^{-4}}\right)^2 \left(\frac{100\ {\rm GeV}}{m_{H_1}}\right)^4 \left(\frac{50\ {\rm GeV}}{m_{\tilde{N}_1}}\right)^2 . \tag{3.38}$$

The upper limits on $\lambda\rho$ are derived by comparing the estimate (3.38) with the current bound on the spin-independent nucleon cross section from the CDMS-II experiment and the expected sensitivities of the SuperCDMS extension [37] and the Xenon1T experiment [38]. These limits are included in Figures 3 and 4. The current bound already excludes large parts of the $(m_{\tilde{N}_1}, \lambda\rho)$-parameter plane, except of the Higgs-boson funnel regions. In the near future, the upgraded experiment SuperCDMS will cover a large part of the parameter space, but it will leave open the lightest Higgs-boson pole region which is theoretically favoured by inflation within the mSUGRA framework. The proposed Xenon1T experiment is expected to further narrow down this uncovered parameter range of the F_D-term model.

Dark Matter may also be indirectly searched for through the detection of its final annihilation products, such as photons, positrons, anti-protons or neutrinos. The dominant channel of the LRHS annihilation in the Higgs funnel is determined by an effective scalar coupling with a $b\bar{b}$ pair, which is approximately independent of the relative velocity of the annihiliating sneutrinos. Rates at low temperatures resulting in gamma-ray or charged particle fluxes are therefore not suppressed compared to the rates at the freeze-out temperature responsible for the LRHS relic density. There are several signals that could be explained as an observation of DM annihilation but, as of now, do not provide a consistent picture interpretable by a single DM candidate and model. For example, the excess in the

diffuse galactic gamma ray spectrum measured by the EGRET detector may be interpreted by a 50-100 GeV WIMP, as given by the LRHS in our model, whereas the 511 keV line observed by the INTEGRAL satellite would hint at an MeV DM particle [39, 40]. Upcoming projects such as the GLAST and PAMELA satellites will have higher sensitivities, probe new energy ranges and should provide a clarification of the observational status. High-energy neutrinos as annihilation products are expected and can be searched for in the Sun and the Earth, as WIMPs can accumulate in their centre. For the LHRS there is no spin-dependent coupling to nuclei, and its capture rate along with the produced neutrino flux is suppressed. In addition, for an annihilation via the Higgs resonance, the effective annihilation coupling required to get the correct relic density is very small. The LRHS is therefore not expected to be within the reach of high-energy neutrino telescopes [41], such as IceCube [42].

4 Conclusions

We have analyzed in detail the relic abundance of the lightest right-handed sneutrinos (LRHS) in the supersymmetric F_D-term model of hybrid inflation. The inflationary potential of the model results from the F-term of the inflaton multiplet $\widehat{S}$. The F_D-term model also includes a subdominant non-anomalous D-term generated from the local $\mathrm{U}(1)_X$ symmetry of the waterfall sector, which does not affect the inflaton dynamics. As was mentioned in the introduction and further discussed in Section 2, the model adequately fits the current CMB data of inflation and provides a natural solution to the so-called gravitino overabundance problem, without resorting to an excessive suppression of possible renormalizable couplings of the inflaton to the MSSM particles. Finally, the F_D-term model closely relates the μ-parameter of the MSSM to an SO(3) symmetric Majorana mass m_N through the VEV of the inflaton field. If $\lambda \sim \rho$, this implies that $\mu \sim m_N$, so the model may naturally predict lepton-number violation at the electroweak scale and potentially account for the BAU via thermal resonant leptogenesis.

In spite of the explicit lepton-number violation through the Majorana term $\frac{1}{2}\,\rho\,\widehat{S}\widehat{N}_i\widehat{N}_i$, the F_D-term hybrid model conserves R-parity. Consequently, the LSP of the spectrum is stable and so qualifies as candidate to address the CDM problem. The new aspect of the F_D-term hybrid model is that thermal right-handed sneutrinos emerge as new candidates to solve this problem, by virtue of the quartic coupling: $\frac{1}{2}\,\lambda\rho\,\widetilde{N}_i^*\widetilde{N}_i^* H_u H_d\ +$ H.c. This new quartic coupling results in the Higgs potential from the F-terms of the inflaton field, and it is not present in the more often-discussed extension of the MSSM, where right-handed

neutrino superfields have bare Majorana masses. Provided that the couplings λ and ρ are not too small, e.g. $\lambda,\, \rho \gtrsim 10^{-2}$, the LRHS $\widetilde{N}_{\rm LSP}$ as LSP can efficiently annihilate via the lightest Higgs-boson resonance H_1 into pairs of b-quarks, in the kinematic region $m_{H_1} \approx 2m_{\widetilde{N}_{\rm LSP}}$, and so drastically reduce its relic density to the observed value: $\Omega_{\rm DM}\, h^2 \approx 0.11$.

Experiments, such as CDMS-II, SuperCDMS and Xenon1T, which look for signatures of WIMPs through their elastic scattering with nuclei, will significantly constrain the allowed parameter space of the F_D-term model. They will exclude most of the parameter space, except possibly of a narrow region close to the lightest H_1-boson resonance, where $m_{H_1} \approx 2m_{\widetilde{N}_{\rm LSP}}$. It might seem that to obtain this particular relation between the masses of the H_1 boson and $\widetilde{N}_{\rm LSP}$, a severe tuning of the model parameters is required. However, it is worth stressing here that such a mass relation may easily be achieved within a mSUGRA framework of the F_D-term model that successfully realizes hybrid inflation.

The LRHS scenario of CDM requires relatively large λ and ρ couplings that could, in principle, make Higgs bosons decay invisibly, e.g. $H \to \widetilde{N}_{\rm LSP}\, \widetilde{N}_{\rm LSP}$. Also, right-handed sneutrinos could be present in the cascade decays of the heavier supersymmetric particles. The collider phenomenology of such a CDM scenario lies beyond the scope of the present article. Instead, we note that the F_D-term hybrid inflationary model can give rise to rich phenomenology which can be probed at high-energy colliders [20, 21], as well as in low-energy experiments of lepton flavour and number violation, such as $0\nu\beta\beta$ decay, $\mu \to e\gamma$ [43], $\mu \to eee$ and $\mu \to e$ conversion in nuclei [44, 45]. It would therefore be very interesting to systematically analyze possible correlations between predictions for cosmological and phenomenological observables in the F_D-term model.

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